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HEAVY CHARGED LEPTON PRODUCTION IN SUPERSTRING INSPIRED E6 MODELS[1,2]

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We investigate the possibility of studying E_6 phenomenology at high energy hadron colliders. The production of heavy leptons pairs *via* a gluon-gluon fusion mechanism is discussed. An enhancement in the parton level cross-section is expected due to the heavy (s)fermion loops which couple to the gluons.

1 Introduction

In this talk we give an overview of a very simple superstring inspired E_6 model followed by an application to heavy charged lepton production, along with some preliminary results and discussion.

1.1 Our Choice of Model

The interest in E_6 as a GUT, from its good old 1970's would-be-"topless" days, was rekindled in late 1984 when Green and Schwarz showed that the $E_8 \otimes E_8'$ gauge group gives rise to anomaly free 10-d string theory. In particular, the Calabi-Yau compactification scheme

$$E_8 \otimes E_8' \;\longrightarrow\; SU(3) \otimes E_6 \otimes E_8' \,,$$

down to 4-d (assuming N=1 SUSY)[3] yields n_g copies of the **27** representation of E_6 , depending on the topology of the compactified space. The SU(3) are the spin connections on the compactified space, and the E_8' couples by gravitational interactions to the matter representations of E_6 ; these interactions should presumably play a role in lifting the supersymmetric mass degeneracy. It should be pointed out that E_6 is by no means a unique GUT group, as other

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dimensional reduction schemes can give rise to different ones. However, a lot of phenomenological work has be done using E_6 and therefore we shall humbly take advantage of this fact.

Further breaking of E_6 , *via* the Hosotani mechanism, leads to rank-6 and rank-5 groups. Here the simplest of these scenarios will be considered, *i.e.*,

$$\mathrm{E}_6 \;\longrightarrow \underbrace{\mathrm{SU(3)_c} \otimes \mathrm{SU(2)_L} \otimes \mathrm{U(1)_Y}}_{\text{Standard Model } (SM)} \otimes \mathrm{U(1)_{Y_E}}\,,$$

where the extra $\mathrm{U(1)_{Y_E}}$ leads to an additional neutral gauge boson, the Z', to the SM. To obtain SM energies the breaking can now proceed by the conventional Higgs mechanism.

The assignment of the quantum numbers to the matter fields in the **27**, with the exception of the SM content, can lead to models which contain leptoquarks, diquarks, or quarks. Here, the assignment that leads to the more familiar SM-like quarks will be chosen, Figure 1. In this model, the third

$$\mathbf{27} = \left\{ \boxed{\begin{pmatrix} u \\ d \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} \nu_e \\ e \end{pmatrix}_{\mathrm{L}} u^c_{\mathrm{L}} d^c_{\mathrm{L}} e^c_{\mathrm{L}}} \; \nu^c_{e_{\mathrm{L}}} d'_{\mathrm{L}} d'^c_{\mathrm{L}} \begin{pmatrix} \nu'_e \\ e' \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} e' \\ \nu'_e \end{pmatrix}^c_{\mathrm{L}} \nu''^c_{e_{\mathrm{L}}} \right\}$$

$$\mathbf{27} = \left\{ \boxed{\begin{pmatrix} c \\ s \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_{\mathrm{L}} c^c_{\mathrm{L}} s^c_{\mathrm{L}} \mu^c_{\mathrm{L}}} \; \nu^c_{\mu_{\mathrm{L}}} s'_{\mathrm{L}} s'^c_{\mathrm{L}} \begin{pmatrix} \nu'_\mu \\ \mu' \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} \mu' \\ \nu'_\mu \end{pmatrix}^c_{\mathrm{L}} \nu''^c_{\mu_{\mathrm{L}}} \right\}$$

$$\mathbf{27} = \left\{ \boxed{\begin{pmatrix} t \\ b \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_{\mathrm{L}} t^c_{\mathrm{L}} b^c_{\mathrm{L}} \tau^c_{\mathrm{L}}} \; \nu^c_{\tau_{\mathrm{L}}} b'_{\mathrm{L}} b'^c_{\mathrm{L}} \begin{pmatrix} \nu'_\tau \\ \tau' \end{pmatrix}_{\mathrm{L}} \begin{pmatrix} \tau' \\ \nu'_\tau \end{pmatrix}^c_{\mathrm{L}} \nu''^c_{\tau_{\mathrm{L}}} \right\}$$

$$Y_E\{\mathbf{27}\} = \left\{ \boxed{\left(\frac{2}{3}\right) \left(\frac{-1}{3}\right) \tfrac{2}{3}\ \tfrac{-1}{3}\ \tfrac{2}{3}} \; \tfrac{5}{3}\ \tfrac{-4}{3}\ \tfrac{-1}{3} \left(\frac{-1}{3}\right) \left(\frac{-4}{3}\right) \ \tfrac{5}{3} \right\}$$

Figure 1: E_6 particle content. On the top three rows, the SM particles are shown in the boxes on the left and their "exotic" counter parts outside the boxes on the right. The exotics have been labeled such that they carry the "expected" SM quantum numbers, with the exception being L=0 for the primed and double primed ones. The bottom row contains the Y_E quantum numbers.

generation sleptons are typically chosen to play the role of the Higgs fields,

$$\Phi_1 = \begin{pmatrix} \tilde{\nu}'_{\tau_{\mathrm{L}}} \\ \tilde{\tau}'_{\mathrm{L}} \end{pmatrix} \equiv \begin{pmatrix} \phi_1^0 \\ \phi_1^- \end{pmatrix}\,,\ \Phi_2 = \begin{pmatrix} \tilde{\tau}'^c_{\mathrm{L}} \\ \tilde{\nu}'^c_{\tau_{\mathrm{L}}} \end{pmatrix} \equiv \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix}\,,\ \Phi_3 = \tilde{\nu}''^c_{\tau_{\mathrm{L}}} \equiv \phi_3^0\,,$$

where $\Phi_k = (\phi^a_{kR} + i\phi^a_{kI})/\sqrt{2}$, with VEV's $\langle \phi^0_{kR} \rangle = v_k$. Therefore, the most general superpotential is of the form

$$W \sim \lambda_1 (u\ d)_{\rm L} i\tau_2 \begin{pmatrix} e' \\ \nu'_e \end{pmatrix}^c_{\rm L} u^c_{\rm L} + \lambda_2 (\nu'_e\ e')_{\rm L} i\tau_2 \begin{pmatrix} u \\ d \end{pmatrix}_{\rm L} d^c_{\rm L} + \lambda_3 (\nu'_e\ e')_{\rm L} i\tau_2 \begin{pmatrix} u \\ d \end{pmatrix}_{\rm L} e^c_{\rm L}$$

$$+\lambda_4 (\nu'_e\ e')^c_{\rm L} i\tau_2 \begin{pmatrix} \nu'_e \\ e' \end{pmatrix}_{\rm L} \nu''^c_{e_{\rm L}} + \lambda_5 d'_{\rm L} d'^c_{\rm L} \nu''^c_{e_{\rm L}} + \dots ,$$

where the generation indices on the Yukawa couplings, λ_a, have been suppressed (*i.e.*, $\lambda_a \sim \lambda^{ijk}_a$ s.t. $i,j,k = 1,2,3$), such that $\lambda^{i33}_4 = \lambda^{i33}_4 = \lambda^{i33}_4 = 0$ for $i = 1,2$. W, and the aforementioned gauge group, specifies all of the couplings in this supersymmetric model,

$$\mathcal{L}_{\rm Yuk} = -\tfrac{1}{2}\left[\tfrac{\partial^2 W}{\partial \tilde{f}_i \partial \tilde{f}_j} f_i f_j + \tfrac{\partial^2 W}{\partial \tilde{f}_i \partial \tilde{f}_j}^{*} \bar{f}_i \bar{f}_j\right] \qquad \underline{\text{Yukawa}}$$

$$V_{\rm Scalar} = \underbrace{\left|\frac{\partial W}{\partial \tilde{f}_i}\right|^2}_{V_F} + \underbrace{\tfrac{1}{2}(g \tilde{f}^*_i T^a_{ij} \tilde{f}_j)^2 + \tfrac{1}{4}(g' Y_i \tilde{f}^*_i \tilde{f}_i)^2 + \tfrac{1}{4}(g'' Y_{E_i} \tilde{f}^*_i \tilde{f}_i)^2}_{V_D} \qquad \underline{\text{Scalar}}$$

where g, g', and g'' ($\approx g'$) are the $\mathrm{SU(2)_L}$, $\mathrm{U(1)_Y}$, and $\mathrm{U(1)_{Y_E}}$ coupling constants, respectively. In order to lift the supersymmetric degeneracy some soft terms are put into the model by hand,

$$V_{\rm Soft} \supseteq \tilde{M}^2_{\rm Q} \left| \begin{pmatrix} \tilde{u}_{\rm L} \\ \tilde{d}_{\rm R} \end{pmatrix} \right|^2 + \tilde{M}^2_u \tilde{u}_{\rm R} \tilde{u}^*_{\rm R} + \dots + \lambda_1 A_u (\tilde{u}\ \tilde{d})_{\rm L} i\tau_2 \underbrace{\begin{pmatrix} \tilde{\tau}'^c_{\rm L} \\ \tilde{\nu}'^c_{\tau_{\rm L}} \end{pmatrix}}_{\text{Higgses}} \tilde{u}^*_{\rm R} + \dots ,$$

with the soft-SUSY breaking terms $\tilde{M}^2_f$ and A_f.

Now, the Higgs potential, V_H ($\subseteq V_{\rm Scalar} + V_{\rm Soft}$), produces masses for all of the particles in the model, with its remaining degrees of freedom creating physical Higgs fields; three neutral-scalars, $H^0_{i=1,2,3}$, a neutral-pseudo-scalar, P^0, and a charged-scalar, $H^\pm$. The mass mixing matrices for the Higgs fields, $\mathcal{M}^2_{\rm Higgs}$, can be obtained from the bilinear terms in V_H,

$$V_H(\phi^a_k) \supseteq \frac{1}{2} \mathcal{M}^2_{ij} (\phi^a_i - \langle \phi^a_i \rangle)(\phi^a_j - \langle \phi^a_j \rangle) \ni \mathcal{M}^2_{ij} = \left. \frac{\partial^2 V_H}{\partial \phi^a_i \partial \phi^a_j} \right|_{VEV's} ,$$

which produces a (2×2) for $\mathcal{M}^2_{H^\pm}$, and two (3×3)'s for $\mathcal{M}^2_{H^0_i}$ and $\mathcal{M}^2_{P^0}$. Diagonalizing these matrices yields the mass terms

$$m^2_{H^\pm} = \frac{\lambda A \nu_3}{\sin(2\beta)} + \left(1 - 2\,\frac{\lambda^2}{g^2}\right) m^2_W , \quad m^2_{P^0} = \frac{\lambda A \nu_3}{\sin(2\beta)} \left(1 + \frac{\nu^2}{4\nu^2_3} \sin^2(2\beta)\right) ,$$

and

$$m^2_{H^0_i} \subseteq (U\mathcal{M}^2_{H^0_i}U^{-1})_D \quad \ni \quad m^2_{H^0_1} \leq m^2_{H^0_2} \leq m^2_{H^0_3}\,,$$

where $\tan\beta = v_2/v_1$, $v^2 = v_1^2 + v_2^2$, $\lambda = \lambda_a^{333}$, and A is a soft term.[b] Assuming a U-gauge, the ϕ_i^a's can be expressed in terms of the physical Higgs fields,

$$\phi_1^{\pm} = \sin\beta\, H^{\pm}\,, \qquad \phi_2^{\pm} = \cos\beta\, H^{\pm}\,,$$

$$\phi_{1I}^0 = \kappa\, v_2 v_3\, P^0\,, \quad \phi_{2I}^0 = \kappa\, v_1 v_3\, P^0\,, \quad \phi_{3I}^0 = \kappa\, v_1 v_2\, P^0\,,$$

and

$$\phi_{iR}^0 = \nu_i + \sum_{j=1}^{3} U_{ij}^{-1} H_j^0\,.$$

where $\kappa = 1/\sqrt{v_1^2 v_2^2 + \nu^2 \nu_3^2}$. This allows for the other masses, and interactions involving the Higgses, to be easily obtained.

The masses for the gauge bosons come from the kinetic terms of the Higgs fields, as follows,

$$\begin{aligned}\mathcal{L}_{\rm K.E.}^{\Phi_i} &\supseteq |(\partial - \tfrac{i}{2}G_\mu)\Phi_i|^2 \\ &\supseteq m_W^2 W_\mu^+ W^{-\mu} + \tfrac{1}{2}(ZZ')_\mu \underbrace{m_Z^2 \begin{pmatrix} 1 & \frac{4v_2^2 - v_1^2}{3v^2/\eta}\sqrt{x_W} \\ \frac{4v_2^2 - v_1^2}{3v^2/\eta}\sqrt{x_W} & \frac{v_1^2 + 16v_2^2 + 25v_3^2}{9v^2/\eta^2}x_W \end{pmatrix}}_{\mathcal{M}^2_{Z\text{-}Z'}} \begin{pmatrix} Z \\ Z' \end{pmatrix}^{\mu},\end{aligned}$$

where

$$G_\mu \supseteq (\tau_3 g\cos\theta_W - g'Y\sin\theta_W)Z_\mu + \sqrt{2}g[\tau_+ W_\mu^- + \tau_- W_\mu^+] + g'' Y_E Z_\mu\,,$$

$m_W = gv/2$, $m_Z = \sqrt{g^2 + g'^2}\, v/2$, and $\eta = g'/g''$. Therefore, the Z and Z' mix to give the mass eigenstates Z_1 and Z_2, with masses $m_{Z_1}(\equiv m_Z) \leq m_{Z_2}$.

The fermion masses come from the Yukawa interaction terms with the Higgs fields,

$$\mathcal{L}_{\rm Yuk} \supseteq -\frac{1}{\sqrt{2}}\left\{\lambda_1\nu_2\,\bar{u}u + \lambda_2\nu_1\,\bar{d}d + \lambda_3\nu_1\,\bar{e}e + \lambda_4\nu_3\,\bar{e}'e' + \lambda_5\nu_3\,\bar{d}'d'\right\}\,,$$

which yields the Yukawa couplings

$$\lambda_1 = \frac{g\, m_u}{\sqrt{2}\, m_W \sin\beta}\,, \quad \lambda_2 = \frac{g\, m_d}{\sqrt{2}\, m_W\cos\beta}\,, \quad \lambda_3 = \frac{g\, m_e}{\sqrt{2}\, m_W\cos\beta}\,,$$

[b] The $\tilde{M}_f^2 \in V_H$ are eliminated by requiring $\partial V_H/\partial\phi_i^a|_{VEV\text{'s}} = 0$.

$$\lambda_4 = \frac{\sqrt{2}}{\nu_3}\, m_{e'}\,, \qquad\qquad \lambda_5 = \frac{\sqrt{2}}{\nu_3}\, m_{d'}\,.$$

The left and right sfermion states, $\tilde{f}_{\rm L,R}$, in general, mix and are obtained from

$$V_{\rm Scalar} + V_{\rm Soft} \supseteq (\tilde{f}_{\rm L}\tilde{f}_{\rm R}) \underbrace{\begin{pmatrix} \mathcal{M}^2_{\rm LL} & \mathcal{M}^2_{\rm LR} \\ \mathcal{M}^2_{\rm LR} & \mathcal{M}^2_{\rm RR} \end{pmatrix}}_{\mathcal{M}^2_{\tilde{f}}} \begin{pmatrix} \tilde{f}^*_{\rm L} \\ \tilde{f}^*_{\rm R} \end{pmatrix} \longrightarrow \sum_{i=1}^{2} m^2_{\tilde{f}_i}\, \tilde{f}_i \tilde{f}^*_i\,,$$

where $\mathcal{M}^2_{\tilde{f}}$ is sfermion mass mixing matrix: e.g. for $\tilde{u}_{\rm L,R}$

$$M^{(\tilde{u})^2}_{LL} = \tilde{M}^2_Q + m^2_u + \tfrac{1}{6}\,(3 - 4x_W)\, m^2_Z\, \cos(2\beta) - \tfrac{1}{36}\, g''^2\, (\nu^2_1 + 4\nu^2_2 - 5\nu^2_3)\,,$$

$$M^{(\tilde{u})^2}_{RR} = \tilde{M}^2_u + m^2_u + \tfrac{2}{3}\, x_W\, m^2_Z\, \cos(2\beta) - \tfrac{1}{36}\, g''^2\, (\nu^2_1 + 4\nu^2_2 - 5\nu^2_3)\,,$$

and

$$M^{(\tilde{u})^2}_{LR} = m_u\, (A_u - m_{e'} \cot\beta)\,.$$

This concludes our survey of a simple E_6 model, as we now have all of the relevant phenomenology for discussing L^+L^- production at hadron colliders.[c]

2 Heavy Lepton Production

An interesting place to look for signatures of new physics is heavy lepton production at high energy hadron colliders, figure 2. For heavy charged leptons,

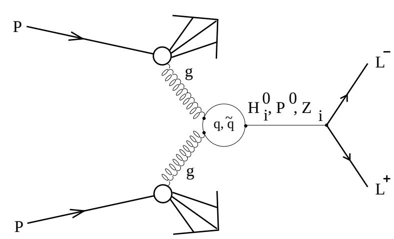

Figure 2: $pp \to gg \to L^+L^-$

$L^\pm$, with mass $m_L \gtrsim \mathcal{O}(150) GeV$ the predominate production mechanism is

[c]For more details on the material covered in this section see [1,3,4,5].

gluon-gluon fusion, within the SM[6] and the minimal supersymmetric standard model ($MSSM$).[7] For high enough energies then, the gluon luminosity of the hadrons can be taken advantage of. Accelerators that may provide good hunting grounds are

- LHC (pp) $14TeV$ $\mathcal{L} \sim 10^5 pb^{-1}/yr\,,$
- TEVATRON $(p\bar{p})$ $1.8TeV$ $\mathcal{L} \sim 10^2 pb^{-1}/yr\,.$

It is expected that the parton level cross-section, $\hat{\sigma}(gg \to L^+L^-)$, should be enhanced due to the number of heavy particles running around in the loop. This process has been computed in the $MSSM$ by Cieza Montalvo, *et al.*,[7] which predicts $\mathcal{O}(10^5)pb^{-1}/yr$ for $50 \leq m_L \leq 400GeV$. Therefore, since E_6 has more particles it is expected that its L^+L^- production rate should be even more enhanced.

2.1 $gg \to L^+L^-$

Figure 3 show the Feynman diagrams that are needed to compute the parton level cross-section (matrix elements) for $gg \to L^+L^-$. These matrix elements

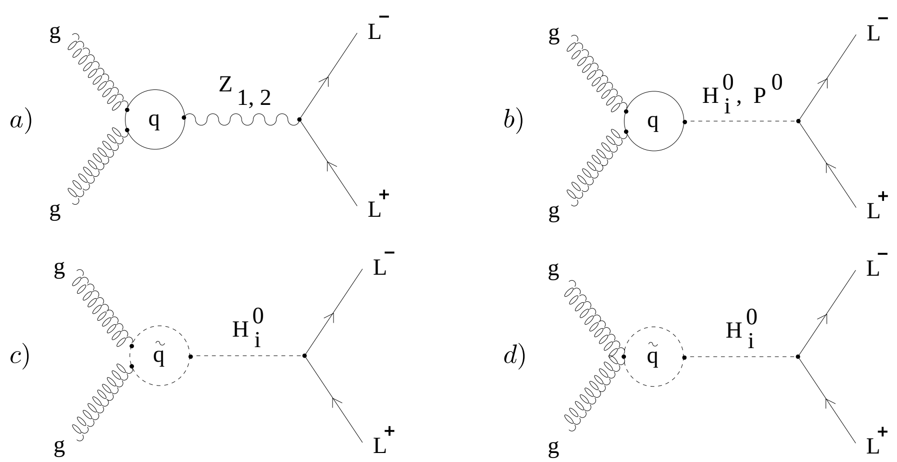

Figure 3: Feynman diagrams for gluon-gluon fusion to heavy charged leptons.

are similar to Cieza Montalvo, *et al.*,[7] and therefore with some care can be

extracted from their paper. The general form of the cross-section is,

$$\hat{\sigma} = \hat{\sigma}_{qZ_{1,2}} + \hat{\sigma}_{qH^0_{1,2,3}} + \hat{\sigma}_{qP^0} + \hat{\sigma}_{\tilde{q}H^0_{1,2,3}} + \underbrace{\hat{\sigma}_{q(Z_{1,2}-P^0)} + \hat{\sigma}_{(\tilde{q}-q)H^0_{1,2,3}}}_{Interference\,Terms} .$$

We have decided not to show the explicit details,[1,2] since they are not very enlightening. However, it worth pointing out that, in the large v_3 limit only the terms involving $Z_{1,2}$ and H_3^0 survive. Before the cross-section can be computed some of the E_6 model parameters need to be constrained.

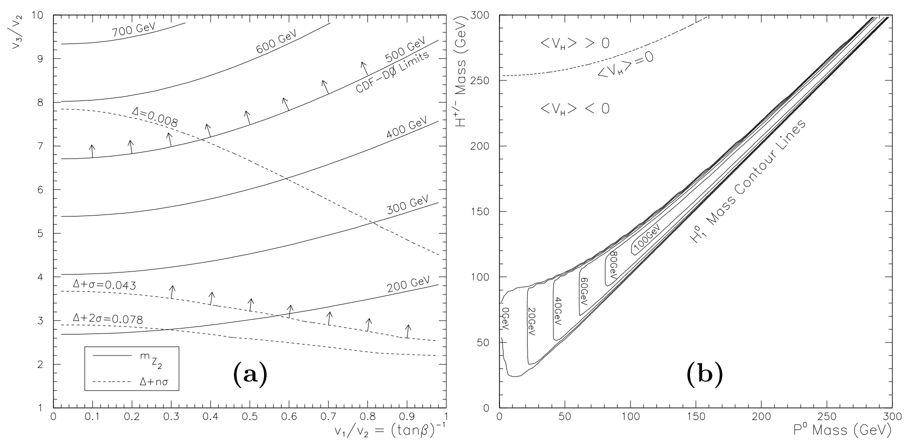

Figure 4: Plots of; (a) m_{Z_2} and Δ contour lines as a function of v_1/v_2 and v_3/v_2. The Δ contour lines are shown at the 0σ, 1σ, and 2σ levels (*cf* [9]). The $m_{Z_2} = 500GeV$ line shows the *CDF* and *DØ* soft-limits.[11] The arrows point to toward the allowed regions on the plot. (b) $m_{H_1^0} \geq 0$ contour lines as a function of m_{P^0} and $m_{H^\pm}$, for $v_1/v_2 = 0.02$ and $v_3/v_2 = 6.7$. The dashed curve in the upper left-hand corner is a plot of the zero of the Higgs potential above which it becomes positive.

The m_{Z_1} (*i.e.*, m_Z) mass can be used to constrain the VEV's by requiring it lies within experimental bounds, *i.e.*,[9]

$$\underbrace{\sin^2\bar{\theta}_W}_{\bar{x}_W} \equiv 1 - \frac{m_W^2}{m_{Z_1}^2} \;<\; \underbrace{\sin^2\theta_W}_{x_W} \equiv \left.\frac{{g'}^2}{g^2 + {g'}^2}\right|_{\mu=m_W} ,$$

due to mixing with m_{Z_2}. Therefore, to within experimental fluctuations in $\bar{x}_W$ ($\approx 0.2247 \pm 0.0019$[8]) and x_W ($\approx 0.233 \pm 0.035|_{\mu=m_W}$) we have

$$\Delta \equiv x_W - \bar{x}_W \approx 0.008 \pm 0.035 .$$

Figure 4.a shows the Δ and m_{Z_2} contour lines as a function of v_1/v_2 (*i.e.*, $1/\tan\beta$) and v_3/v_2. Here, $v_1/v_2 \lesssim \mathcal{O}(1)$ has been chosen, since $m_b << m_t$ for any reasonable range of Yukawa couplings.[9,10] Taking $\Delta + 1\sigma$ contour line gives $v_3/v_2 \gtrsim \mathcal{O}(3.5)$ (cf [3,9]) or $m_{Z_2} \gtrsim \mathcal{O}(200) GeV$. Since the error in Δ is fairly large, a more conservative approach has been taken by invoking the CDF and $D\emptyset$ soft-limits (*i.e.*, assuming SM couplings) on m_{Z_2} (fig. 4.a).[11] These give fairly reasonable bounds on m_{Z_2}, *i.e.*, $m_{Z_2} \gtrsim \mathcal{O}(500) GeV$ ($\Rightarrow v_3/v_2 \gtrsim \mathcal{O}(7.5)$), since Y_E's $\sim \mathcal{O}(Y)$'s (cf fig. 1). For the moment let us make the further assertion that $\lambda_1 \sim \lambda_2$. Then we have $v_1/v_2 = 0.02$ (*i.e.*, $\sim \mathcal{O}(m_b/m_t)$), which implies $v_3/v_2 = 6.7$ for $m_{Z_2} \approx \mathcal{O}(500) GeV$.

Now that the VEV's have been fixed, the Higgses masses can be constrained by adjusting λ and A, or equivalently m_{P^0} and $m_{H^\pm}$. Figure 4.b shows the contour lines for the lightest scalar-Higgs mass, $m_{H_1^0}$, as a function of m_{P^0} and $m_{H^\pm}$. Notice that variation in m_{P^0} and $m_{H^\pm}$ is restricted to a very narrow region within which $m_{H_1^0} \geq 0$. Also, for variations of $m_{P^0}, m_{H^\pm} \lesssim \mathcal{O}(1) TeV$ it turns out that $m_{H_3^0}$ remains fairly constant. Therefore, we are free to choose m_{P^0} and $m_{H^\pm}$ as we like since L^+L^- production becomes insensitive to $m_{H_{1,2}^0}$ and m_{P^0} in the large v_3 limit (which is the case here). Furthermore, the region over which m_{P^0} and $m_{H^\pm}$ are allowed to vary changes insignificantly for $v_1/v_2 \lesssim \mathcal{O}(1)$, which allows us to recant our previous assertion. Here we will consider $m_{P^0} = 200 GeV$, $m_{H^\pm} \approx \mathcal{O}(214) GeV$, $6.7 \leq v_3/v_2 \leq 9.1$, and $0.02 \leq v_1/v_2 \leq 0.9$.

In general, the soft-terms should be evolved down from some SUSY unification scale to give proper masses to the scalar squarks. However, as the details of this evolution have not been completely settled, it is typical to treat these terms as parameters. Here we will choose $\tilde{M}_{\tilde{f}} = A_f \equiv m_S$, such that $\mathcal{O}(0.4) TeV \lesssim m_S \lesssim \mathcal{O}(1) TeV$.

The exotic quark masses, the $m_{q'}$'s, will be assumed degenerate, such that $\mathcal{O}(200) GeV \lesssim m_{q'} \lesssim \mathcal{O}(600) GeV$.

Finally, the $e'^{\pm}$ will be designated to play the role of the heavy charged lepton, $L^\pm$.

2.2 Results

Figure 5.a shows the rapidity distribution, $\partial\sigma/\partial y$, at rapidity zero, $y = 0$, as a function of heavy lepton mass, m_L, for a typical set of E_6 parameters at LHC.[d] Notice that the terms involving P^0 are suppressed, as advertised. The terms

[d] The rapidity distribution was obtained by folding the parton level cross-section in with the $DO1.1$ gluon distribution function.[12]

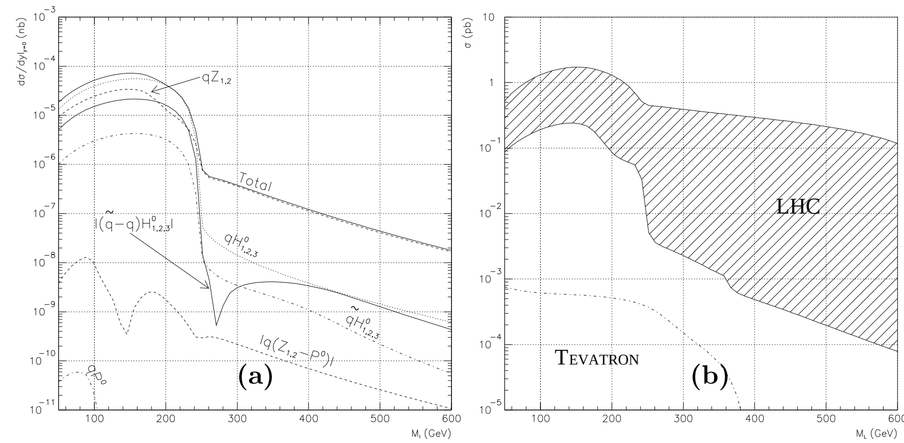

Figure 5: Plots of; (a) the rapidity distribution (in pb) at $y = 0$ for heavy charged lepton production at LHC as a function of heavy lepton mass, where $v_1/v_2 = 0.02$, $v_3/v_2 = 6.7$, and $m_S = 400 GeV$. The mass spectrum for the non-SM particles are, $m_{Z_2} \approx 496 GeV$ ($\Gamma_{Z_2} \approx 20.9 GeV$), $m_{P^0} \approx 200 GeV$ ($\Gamma_{P^0} \approx 16.4 GeV$), $m_{H^\pm} \approx 215 GeV$, $m_{H_1^0} \approx 94.3 GeV$ ($\Gamma_{H_1^0} \approx 7.50 \times 10^{-3} GeV$), $m_{H_2^0} \approx 200 GeV$ ($\Gamma_{H_2^0} \approx 16.5 GeV$), $m_{H_3^0} \approx 495 GeV$ ($\Gamma_{H_3^0} \approx 0.230 GeV$), $m_{q'} = 200 GeV$. (b) the results for the total L^+L^- production cross-section (in nb) at LHC and the TEVATRON as a function of m_L. The hashed region are the LHC results.

involving $H^0_{1,2}$ are also suppressed, but are not explicitly shown. The dramatic drop in $\partial\sigma/\partial y|_{y=0}$ at $m_L \approx 250 GeV$ corresponds the Z_2 and H_3^0 resonance cut-off's. The cut-off's vary with the Z_2 mass, since $m_{H_3^0} \approx m_{Z_2} \approx \mathcal{O}(g''^2 v_3^2)$ for large v_3 (cf[10]). Not so obvious in this plot, is the slight kink in the $qZ_{1,2}$ curve around $m_L \approx 200 GeV$ (*i.e.*, $\approx m_{q'}$). This corresponds to the heavy quarks in the loops going off shell. The effect becomes more noticable as the exotic quark masses are pushed up to $\mathcal{O}(600) GeV$, because it sustains $\partial\sigma/\partial y|_{y=0}$ at $\mathcal{O}(10^{-4}) pb$ before gradually starting to drop off.

Finally, Figure 5.b shows the total L^+L^- production cross-section at LHC over the E_6 model parameter space specified in § 2.1. Also shown are the results for the TEVATRON, with the parameters given in figure 5.a . Clearly not enough events are produced at the TEVATRON to make a search worthwhile, *i.e.*, $\lesssim \mathcal{O}(0.1) events/yr$. However, at LHC we expect $\mathcal{O}(10^{4\pm1}) events/yr$. Unfortunately, this a factor of at least 10 less than predicted by the $MSSM$.

The reason for this discrepancy is because v_3 is constrained to be large. In this limit the L^+L^- production losses out over $MSSM$ since the number of Higgs propagators has been reduce from four to one, whereas $MSSM$ has

three. Indeed, when v_3/v_2 drops below the $\Delta + 2\sigma$ contour line, in figure 4.a, the other Higgses start to contribute and L^+L^- production increases by a factor of $\mathcal{O}(10)$, for $m_L \lesssim \mathcal{O}(100)GeV$.

It should also be pointed that, for $MSSM$ model parameters which yield a particle spectrum similar (*i.e.*, with comparable masses) to that of E_6, a fair chunk of its parameter space is eliminated by unitarity constraints.[7] In particular, for the results shown in figure 5.a, $MSSM$ production is restricted to the region $m_L \lesssim \mathcal{O}(250)GeV$, for $m_{H_1^0} \gtrsim \mathcal{O}(600)GeV$ and $\tan\beta \lesssim \mathcal{O}(5)$. For more conservative $MSSM$ parameters, $m_L \lesssim \mathcal{O}(400)GeV$.

3 Closing Remarks

A simple E_6 model was constructed and used to compute L^+L^- production at high energy hadron colliders. We expect $\mathcal{O}(10^{4\pm1})events/yr$ at LHC and "zero" at the TEVATRON, for $50GeV \leq m_L \leq 600GeV$. The results were a factor of at least 10 less than the $MSSM$ results due to the CDF and $D\emptyset$ soft-limits on m_{Z_2},[11] which caused the $H^0_{1,2}$ and P^0 contributions to become suppressed.

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